Neutrino Mass and Dark Matter from Gauged B−L Breaking *

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We discuss a new radiative seesaw model with the gauged B−L symmetry which is spontaneously broken. We improve the previous model by using the anomaly-free condition without introducing too many fermions. In our model, dark matter, tiny neutrino masses and neutrino oscillation data can be explained simultaneously, assuming the B−L symmetry breaking at the TeV scale.

I. INTRODUCTION

The neutrino oscillation data [2–4] have shown us neutrinos have tiny masses. If ν_R are introduced to the standard model of particle physics (SM), there are two possible mass terms for neutrinos (See e.g., Ref. [5]), the Dirac type $\overline{\nu_L}\nu_R$ and the Majorana type $\overline{(\nu_R)^c}\nu_R$. In radiative seesaw models (See e.g., Refs. [6–12]), an *ad hoc* unbroken Z_2 symmetry forbids generating neutrino masses at the tree level and explains the dark matter (DM) stability. A model in Ref. [12] was constructed such that the breaking of the $\mathrm{U(1)_{B-L}}$ gauge symmetry gives a residual symmetry for the DM stability and the Majorana neutrino mass of ν_R. However, the anomaly cancelation for the $\mathrm{U(1)_{B-L}}$ gauge symmetry requires to introduce more additional fermions except for particles for the radiative neutrino mass.

In this talk, we propose a new model which is an improved version of the model in Ref. [12] from the view point of the anomaly cancellation. With appropriate $\mathrm{U(1)_{B-L}}$ charge assignments, there exists an unbroken global U(1) symmetry even after the breakdown of the $\mathrm{U(1)_{B-L}}$ symmetry. The global U(1) symmetry stabilizes the DM, so that we hereafter call it $\mathrm{U(1)_{DM}}$. In our work, the DM candidate is a new scalar boson. Furthermore, the Dirac mass term of neutrinos is radiatively generated at the one-loop level due to the quantum effect of the new particles. Tiny neutrino masses are explained by the two-loop diagrams with a Type-I-Seesaw-like mechanism. We find that the model can satisfy current data from the neutrino oscillation, the lepton flavor violation (LFV), the relic abundance and the direct search for the DM, and the LHC experiment.

II. MODEL

We introduce new particles which listed in Table I. We determine assignment of $\mathrm{U(1)_{B-L}}$ charges from conditions for cancellation of the $[\mathrm{U(1)_{B-L}}] \times [\text{gravity}]^2$ and $[\mathrm{U(1)_{B-L}}]^3$ anomalies;

$$3-\frac{1}{3}N_{\nu_R}-\frac{2}{3}N_\psi=0, \quad 3-\frac{1}{27}N_{\nu_R}+\left(-2x^2-\frac{4}{3}x-\frac{8}{27}\right)N_\psi=0, \tag{1}$$

where N_ψ is the number of ψ_{Ri} (the same as the number of ψ_{Li}), and N_{ν_R} is the number of ν_{Ra}.

There are four solutions as presented in Table II. Except for Case III, the $\mathrm{U(1)_{B-L}}$ charges of some new particles are irrational numbers while the $\mathrm{U(1)_{B-L}}$ symmetry is spontaneously broken by the vacuum expectation value (VEV) of σ^0 whose $\mathrm{U(1)_{B-L}}$ charge is a rational number. Therefore, the irrational charges are conserved, and the lightest particle with an irrational $\mathrm{U(1)_{B-L}}$ charge becomes stable so that the particle can be regarded as a DM candidate. In this talk, we take Case IV as an example.

In addition to the SM one, the new Yukawa interactions are given by

$$\mathcal{L}_\mathrm{Y}=-(y_R)_i\,\overline{(\nu_R)_i}\,(\nu_R)^c_i\,(\sigma^0)^*-(y_\psi)_i\,\overline{(\psi_R)_i}\,(\psi_L)_i\,(\sigma^0)^*-h_{ij}\,\overline{(\psi_L)_i}\,(\nu_R)_j\,s^0-f_{\ell i}\,\overline{(L_L)_\ell}\,(\psi_R)_i\,\tilde{\eta}+\text{h.c.}, \tag{2}$$

where $\tilde{\eta}\equiv i\sigma_2\,\eta^*$. The scalar potential in our model is the same as that in the previous model [12]:

$$V=-\mu_\phi^2\Phi^\dagger\Phi+\mu_s^2|s^0|^2+\mu_\eta^2\eta^\dagger\eta-\mu_\sigma^2|\sigma^0|^2+\mu_3\left(s^0\,\eta^\dagger\,\Phi+\text{h.c.}\right)+\lambda_\phi\left(\Phi^\dagger\Phi\right)^2+\lambda_s|s^0|^4+\lambda_\eta\left(\eta^\dagger\eta\right)^2+\lambda_\sigma|\sigma^0|^4$$
$$+\lambda_{s\sigma}|s^0|^2|\sigma^0|^2+\lambda_{s\eta}|s^0|^2\eta^\dagger\eta+\lambda_{s\phi}|s^0|^2\Phi^\dagger\Phi+\lambda_{\sigma\eta}|\sigma^0|^2\eta^\dagger\eta+\lambda_{\sigma\phi}|\sigma^0|^2\Phi^\dagger\Phi+\lambda_{\phi\phi}(\eta^\dagger\eta)(\Phi^\dagger\Phi)+\lambda_{\eta\phi}(\eta^\dagger\Phi)(\Phi^\dagger\eta). \tag{3}$$

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* This talk is based on Ref. [1].

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TABLE I: Particle contents in this model. Indices i and a run from 1 to N_ψ and from 1 to N_{ν_R}, respectively.

	σ^0	$(\nu_R)_a$	$(\psi_L)_i$	$(\psi_R)_i$	η	s^0
$\mathrm{SU(2)_I}$	**1**	**1**	**1**	**1**	**2**	**1**
$\mathrm{U(1)_Y}$	0	0	0	0	1/2	0
$\mathrm{U(1)}_{B-L}$	2/3	−1/3	$x+2/3$	x	$x+1$	$x+1$
Spin	0	1/2	1/2	1/2	0	0

TABLE II: Sets of N_ψ, N_{ν_R} and x, for which the $\mathrm{U(1)_{B-L}}$ gauge symmetry is free from anomaly.

	Case I	Case II	Case III	Case IV
N_ψ	1	2	3	4
N_{ν_R}	7	5	3	1
x	$\frac{2\sqrt{3}-1}{3}$	$\frac{\sqrt{6}-1}{3}$	$\frac{1}{3}$	$\frac{\sqrt{3}-1}{3}$

Neutral scalar fields are given by $\phi^0 = \frac{1}{\sqrt{2}}(\phi^0_r + i z_\phi), \sigma^0 = \frac{1}{\sqrt{2}}(\sigma^0_r + i z_\sigma), \eta^0 = \frac{1}{\sqrt{2}}(\eta^0_r + i\eta^0_i), s^0 = \frac{1}{\sqrt{2}}(s^0_r + i s^0_i)$. Two scalar fields ϕ^0 and σ^0 obtain VEVs v_ϕ $[= \sqrt{2}\,\langle\phi^0\rangle = 246\,\mathrm{GeV}]$ and v_σ $[= \sqrt{2}\,\langle\sigma^0\rangle]$. The VEV v_σ provides a mass of the $\mathrm{U(1)_{B-L}}$ gauge boson Z' as $m_{Z'} = (2/3)g_{\mathrm{B-L}}v_\sigma$, where $g_{\mathrm{B-L}}$ is the $\mathrm{U(1)_{B-L}}$ gauge coupling constant. After the gauge symmetry breaking with v_ϕ and v_σ, we can confirm in Eqs. (2) and (3) that there is a residual global $\mathrm{U(1)_{DM}}$ symmetry, for which irrational $\mathrm{U(1)_{B-L}}$-charged particles (η, s^0, ψ_{Li}, and ψ_{Ri}) have the same $\mathrm{U(1)_{DM}}$-charge while the other particles are neutral.

Two CP-even scalar particles h^0 and H^0 are obtained by ϕ^0-σ^0 mixing as $\sin 2\theta_0 = \frac{2\lambda_{\sigma\phi}v_\phi v_\sigma}{m^2_{H^0} - m^2_{h^0}}$. Two neutral complex scalars η^0 and s^0 are obtained by η^0-s^0 mixing as $\sin 2\theta'_0 = \frac{\sqrt{2}\mu_3 v_\phi}{m^2_{\mathcal{H}^0_2} - m^2_{\mathcal{H}^0_1}}$. Scalar masses are given by

$$m^2_{h^0,H^0} = \lambda_\phi v^2_\phi + \lambda_\sigma v^2_\sigma \mp \sqrt{\left(\lambda_\phi v^2_\phi - \lambda_\sigma v^2_\sigma\right)^2 + \lambda^2_{\sigma\phi} v^2_\phi v^2_\sigma},\ m^2_{\mathcal{H}^0_1,\mathcal{H}^0_2} = \frac{1}{2}\left(m^2_\eta + m^2_s \mp \sqrt{\left(m^2_\eta - m^2_s\right)^2 + 2\mu^2_3 v^2_\phi}\right), \tag{4}$$

where $m^2_\eta = \mu^2_\eta + (\lambda_{\phi\phi} + \lambda_{\eta\phi})\, v^2_\phi/2 + \lambda_{\sigma\eta} v^2_\sigma/2$, $m^2_s = \mu^2_s + \lambda_{s\phi} v^2_\phi/2 + \lambda_{s\sigma} v^2_\sigma/2$. The mass of the charged scalar $\eta^\pm$ is $m^2_{\eta^\pm} = m^2_\eta - \lambda_{\eta\phi}\, v^2_\sigma/2$. Nambu-Goldstone bosons z_ϕ and z_σ are absorbed by Z and Z' bosons, respectively.

III. PHENOMENOLOGY

A. Neutrino masses

Tiny neutrino masses are generated by two-loop diagrams in Fig. 1 [12]. The mass matrix m_ν is expressed in the flavor basis as

$$(m_\nu)_{\ell\ell'} \;=\; \sum_{i,j,a} f_{\ell i}\, h_{ia}\, (m_R)_a\, (h^T)_{aj}\, (f^T)_{j\ell'} \Big[(I_1)_{ija} + (I_2)_{ija}\Big]/(16\pi^2)^2, \tag{5}$$

where explicit formulas of $(I_1)_{ija}$ and $(I_2)_{ija}$ are shown in Ref. [1]. The neutrino mass matrix $(m_\nu)_{\ell\ell'}$ is diagonalized by a unitary matrix U_{MNS}, the so-called Maki-Nakagawa-Sakata (MNS) matrix [13], as $U^\dagger_{\mathrm{MNS}}\, m_\nu\, U^*_{\mathrm{MNS}} = \mathrm{diag}(m_1 e^{i\alpha_1},\, m_2 e^{i\alpha_2},\, m_3 e^{i\alpha_3})$. We take m_i (i = 1-3) to be real and positive values. Two differences of three phases α_i are physical Majorana phases. In our analysis, the following values [2–4] obtained by neutrino oscillation measurements are used in order to search for a benchmark point of model parameters:

$$\sin^2 2\theta_{23} = 1, \quad \sin^2 2\theta_{13} = 0.09, \quad \tan^2\theta_{12} = 0.427, \quad \delta = 0, \quad \big\{\alpha_1,\, \alpha_2,\, \alpha_3\big\} = \big\{0,\, 0,\, 0\big\}, \tag{6}$$

$$m_1 = 10^{-4}\,\mathrm{eV}, \quad \Delta m^2_{21} = 7.46\times 10^{-5}\,\mathrm{eV}^2, \quad \Delta m^2_{32} = +2.51\times 10^{-3}\,\mathrm{eV}^2, \quad \text{where } \Delta m^2_{ij} \equiv m^2_i - m^2_j. \tag{7}$$

By using an ansatz [1] for the structure of Yukawa matrix $f_{\ell i}$, we found a benchmark point as

$$f = \begin{pmatrix} 1.79 & -2.49 & -1.97 & 2.56 \\ -1.82 & 1.10 & 1.30 & -0.818 \\ 1.40 & -0.598 & -0.905 & 0.222 \end{pmatrix} \times 10^{-2},\ h = \begin{pmatrix} 0.7 & 0.8 & 0.9 & 1 \end{pmatrix}^T,\ \big\{g_{\mathrm{B-L}},\, m_{Z'}\big\} = \big\{0.1,\, 4\,\mathrm{TeV}\big\}, \tag{8}$$

$$\big\{m_{h^0},\, m_{H^0},\, \cos\theta_0\big\} = \big\{125\,\mathrm{GeV},\, 1\,\mathrm{TeV},\, 1\big\},\ \big\{m_{\mathcal{H}^0_1},\, m_{\mathcal{H}^0_2},\, \cos\theta'_0\big\} = \big\{60\,\mathrm{GeV},\, 450\,\mathrm{GeV},\, 0.05\big\}, \tag{9}$$

$$m_{\eta^\pm} = 420\,\mathrm{GeV},\ (m_R)_1 = 250\,\mathrm{GeV},\ \big\{m_{\psi_1},\, m_{\psi_2},\, m_{\psi_3},\, m_{\psi_4}\big\} = \big\{650\,\mathrm{GeV},\, 750\,\mathrm{GeV},\, 850\,\mathrm{GeV},\, 950\,\mathrm{GeV}\big\}. \tag{10}$$

The values of $\{m_{h^0},\, m_{H^0},\, \cos\theta_0\}$ correspond to $\lambda_\phi \simeq 0.13$, $\lambda_\sigma \simeq 2.8\times 10^{-4}$ and $\lambda_{\sigma\phi} = 0$. The values of $\{m_{\mathcal{H}^0_1},\, m_{\mathcal{H}^0_2},\, \cos\theta'_0\}$ and $m_{\eta^\pm}$ can be produced by $m_s \simeq 60\,\mathrm{GeV}$, $m_\eta \simeq 450\,\mathrm{GeV}$, $\mu_3 \simeq 57\,\mathrm{GeV}$ and $\lambda_{\eta\phi} \simeq 0.86$.

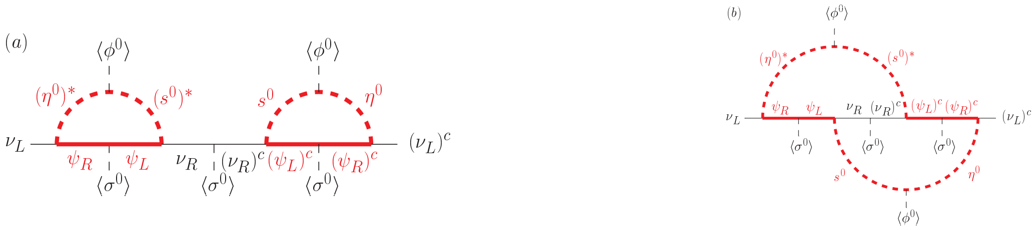

FIG. 1: Two-loop diagrams for tiny neutrino masses in this model.

B. Lepton flavor violation

We consider the condition of the LFV decays of charged leptons. The charged scalar $\eta^\pm$ contributes to the branching ratio (BR) of $\mu \to e\gamma$ whose formula have been calculated [14]. At the benchmark point, we have $\mathrm{BR}(\mu \to e\gamma) = 6.1 \times 10^{-14}$ which satisfies the current constraint $\mathrm{BR}(\mu \to e\gamma) < 5.7 \times 10^{-13}$ (90% C.L.) [15].

C. Dark matter

In our model, the scalar $\mathcal{H}_1^0$ turns out to be the DM candidate due to the following reason. If the DM is the fermion ψ_1, it annihilates into a pair of SM particles via the s-channel process mediated by h^0 and H^0. Even for a maximal mixing $\cos\theta_0 = 1/\sqrt{2}$ [16], the observed abundance of the DM [17] requires $v_\sigma \lesssim 10\,\mathrm{TeV}$. The current constraint from direct searches of the DM [18] requires larger v_σ in order to suppress the Z' contribution.

The scalar DM $\mathcal{H}_1^0$ at the benchmark point is dominantly made from s^0 which is a gauge-singlet field under the SM gauge group, because of the tiny mixing $\cos\theta_0' = 0.05$. The annihilation of $\mathcal{H}_1^0$ into a pair of the SM particles is dominantly caused by the s-channel scalar mediation via h^0 [19] because H^0 is assumed to be heavy. The coupling constant $\lambda_{\mathcal{H}_1^0\mathcal{H}_1^0 h^0}$ for the $\lambda_{\mathcal{H}_1^0\mathcal{H}_1^0 h^0}\, v_\phi \mathcal{H}_1^0 \mathcal{H}_1^{0*} h^0$ interaction controls the annihilation cross section, the invisible decay $h^0 \to \mathcal{H}_1^0 \mathcal{H}_1^{0*}$ in the case of kinematically accessible, and the h^0 contribution to the spin-independent scattering cross section σ_{SI} on a nucleon. In Ref. [20], for example, we see that $\mathcal{H}_1^0$ with $m_{\mathcal{H}_1^0} = 60\,\mathrm{GeV}$ and $\lambda_{\mathcal{H}_1^0\mathcal{H}_1^0 h^0} \sim 10^{-3}$ can satisfy constraints from the relic abundance of the DM and the invisible decay of h^0. We see also that the h^0 contribution to σ_{SI} is small enough to satisfy the current constraint $\sigma_{\mathrm{SI}} < 9.2 \times 10^{-46}\,\mathrm{cm}^2$ for $m_{\mathrm{DM}} = 60\,\mathrm{GeV}$ [18]. Although the scattering of $\mathcal{H}_1^0$ on a nucleon is mediated also by the Z' boson in this model, the contribution can be suppressed by taking a large v_σ. The benchmark point corresponds to $v_\sigma = 60\,\mathrm{TeV}$ and gives about $6.6 \times 10^{-47}\,\mathrm{cm}^2$ for the scattering cross section via Z', which is smaller than the current constraint [18] by an order of magnitude. Thus, the constraint from the direct search of the DM is also satisfied at the benchmark point.

D. Z' and ν_R search

The LEP-II bound $m_{Z'}/g_{\mathrm{B-L}} \gtrsim 7\,\mathrm{TeV}$ [21] is satisfied at the benchmark point because of $m_{Z'}/g_{\mathrm{B-L}} = 40\,\mathrm{TeV}$ which we take for a sufficient suppression of σ_{SI} for the direct search of the DM. The production cross section of Z' with $g_{\mathrm{B-L}} = 0.1$ and $m_{Z'} = 4\,\mathrm{TeV}$ is about 0.3 fb at the LHC with $\sqrt{s} = 14\,\mathrm{TeV}$ [22]. Notice that the current bound $m_{Z'} \gtrsim 3\,\mathrm{TeV}$ at the LHC [23] is for the case where the gauge coupling for Z' is the same as the one for Z, namely $g_{\mathrm{B-L}} \simeq 0.7$. Decay branching ratios of Z' are shown at the benchmark point in Table III. Decays of ψ_i are dominated by $\psi_i \to \nu_R \mathcal{H}_1^0$ with the Yukawa coupling constants h_{i1} because $y_{\ell i}$ for $\psi_i \to \ell^\pm \eta^\mp$ are small in order to satisfy the $\mu \to e\gamma$ constraint. The $\mathcal{H}_2^0$ ($\simeq \eta^0$) decays into $h^0 \mathcal{H}_1^0$ via the trilinear coupling constant μ_3. The main decay mode of $\eta^\pm$ is $\eta^\pm \to W^\pm \mathcal{H}_1^0$ through the mixing θ_0' between η^0 and s^0.

The ν_R decay into H^0 is forbidden because it is heavier than ν_R at the benchmark point. Since the B$-$L charge of ν_R is rather small, ν_R is not produced directly from Z'. However, ν_R can be produced through the decays of ψ_i. As a result, about 18 % of Z' produces ν_R. For $\nu_R \to W\ell$ (56 %) followed by the hadronic decay of W (68 %), the ν_R would be reconstructed. In this model, an invariant mass of a pair of the reconstructed ν_R is not at $m_{Z'}$ in contrast with a naive model where only three ν_R with B$-$L $= -1$ are introduced to the SM. This feature of ν_R also enables us to distinguish this model from the previous model in Ref. [12] where ν_R with B$-$L $= 1$ can be directly produced by the Z' decay.

TABLE III: Branching ratios of Z' decays.

$q\,\bar{q}$	$\ell\,\bar{\ell}$	$\nu_L\overline{\nu_L}$	$\nu_R\overline{\nu_R}$	$\psi_1\overline{\psi_1}$	$\psi_2\overline{\psi_2}$	$\psi_3\overline{\psi_3}$	$\psi_4\overline{\psi_4}$	$\mathcal{H}_1^0\mathcal{H}_1^{0*}$	$\mathcal{H}_2^0\mathcal{H}_2^{0*}$	$\eta^+\eta^-$
0.21	0.32	0.16	0.0059	0.046	0.045	0.044	0.043	0.041	0.038	0.039

IV. CONCLUSIONS

We have improved the model in Ref. [12] by considering anomaly cancellation of the $\mathrm{U}(1)_{\mathrm{B-L}}$ gauge symmetry. We have shown that there are four anomaly-free cases of B−L charge assignment, and three of them have an unbroken global $\mathrm{U}(1)_{\mathrm{DM}}$ symmetry. The $\mathrm{U}(1)_{\mathrm{DM}}$ guarantees that the lightest $\mathrm{U}(1)_{\mathrm{DM}}$-charged particle is stable such that it can be regarded as a DM candidate. The spontaneous breaking of the $\mathrm{U}(1)_{\mathrm{B-L}}$ symmetry generates the Majorana mass term of ν_R and masses of new fermions ψ. In addition, the Dirac mass term of neutrinos is generated at the one-loop level where the DM candidate involved in the loop. Tiny neutrino masses are obtained at the two-loop level.

The case of the fermion DM is excluded, and the lightest $\mathrm{U}(1)_{\mathrm{DM}}$-charged scalar $\mathcal{H}_1^0$ should be the DM in this model. We have found a benchmark point of model parameters which satisfies current constraints from neutrino oscillation data, lepton flavor violation searches, the relic abundance of the DM, direct searches for the DM, and the LHC experiments. In such radiative seesaw models, ν_R would be produced at the LHC. In our model, ν_R cannot be directly produced by the Z' decay, but can be produced by the cascade decay $Z' \to \psi_i\overline{\psi}_i \to \nu_R\overline{\nu_R}\mathcal{H}_1^0\mathcal{H}_1^{0*}$. By the unusual B−L charge of ν_R, the invariant mass distribution of $\nu_R\overline{\nu_R}$ does not take a peak at $m_{Z'}$, which could be a characteristic signal.

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